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Phenomenology of Mixed Modulus-Anomaly Mediation in Fluxed String Compactifications and Brane Models

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Abstract: In some string compactifications, for instance the recently proposed KKLT set-up, light moduli are stabilized by nonperturbative effects at supersymmetric AdS vacuum which is lifted to a dS vacuum by supersymmetry breaking uplifting potential. In such models, soft supersymmetry breaking terms are determined by a specific mixed modulus-anomaly mediation in which the two mediations typically give comparable contributions to soft parameters. Similar pattern of soft terms can arise also in brane models to stabilize the radion by nonperturbative effects. We examine some phenomenological consequences of this mixed modulus-anomaly mediation, including the pattern of low energy sparticle spectrum and the possibility of electroweak symmetry breaking. It is noted that adding the anomaly-mediated contributions at M_{GUT} amounts to replacing the messenger scale of the modulus mediation by a mirage messenger scale $(m_{3/2}/M_{Pl})^{\alpha/2}M_{GUT}$ where $\alpha = m_{3/2}/[M_0 \ln(M_{Pl}/m_{3/2})]$ for M_0 denoting the modulus-mediated contribution to the gaugino mass at M_{GUT}. The minimal KKLT set-up predicts $\alpha = 1$. As a consequence, for $\alpha = \mathcal{O}(1)$, the model can lead to a highly distinctive pattern of sparticle masses at TeV scale, particularly when $\alpha = 2$.

Keywords: Supergravity Models, Supersymmetry Breaking, Supersymmetric Standard Model.

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1. Introduction

Low energy supersymmetry (SUSY) is one of the prime candidates for physics beyond the standard model at TeV scale [1]. One of the central questions in supersymmetric models is to understand the mechanism of SUSY breaking, in particular the origin of soft SUSY breaking terms in the low energy effective lagrangian [2]. Most phenomenological aspects of supersymmetric models are determined by those soft terms which would be induced by the auxiliary components of SUSY breaking messenger fields. In string theory, the most plausible candidates for messenger fields are the moduli fields (including dilaton) describing the continuous degeneracy of string vacua at leading approximation [3]. In addition to string moduli, the 4-dimensional supergravity (SUGRA) multiplet provides a model-independent source of SUSY breaking, i.e. the anomaly-mediation [4], which generically induces the soft masses $m_{soft} \sim m_{3/2}/8\pi^2$ where $m_{3/2}$ is the gravitino mass. To identify the dominant source of soft terms, one needs to understand how those moduli are stabilized at a nearly 4D Poincare invariant vacuum. For instance, if some moduli ϕ_H are stabilized with a heavy mass $m_{\phi_H} \gg 8\pi^2 m_{3/2}$, those moduli would have the auxiliary component $F_{\phi_H} \sim m_{3/2}^2/m_{\phi_H} \ll m_{3/2}/8\pi^2$ (in the unit with $M_{Pl} = 1$), thus their contribution to soft terms are negligible even compared to the anomaly-mediated ones. On the other hand, light moduli ϕ_L with a mass $m_{\phi_L} \sim m_{3/2}$ can have $F_{\phi_L} \sim m_{3/2}$ which might provide $m_{soft} \sim m_{3/2}$ dominating over the anomaly-mediation. In addition to stabilizing all the relevant moduli, one needs to make sure that the resulting vacuum energy density is tuned to be nearly zero since the soft scalar masses can be affected by any additional source of the vacuum energy density [5].

Recently KKLT has proposed an interesting set-up to stabilize the moduli within the framework of type IIB string theory. The KKLT set-up [6] involves three steps to achieve a SUSY breaking Minkowski (or de Sitter) vacuum, while stabilizing all (or most of) moduli. The first step is to introduce the NS and RR 3-form fluxes, H_3 and F_3, stabilizing the dilaton S and all complex structure moduli Z_α. For some flux vacua, $G_3 = H_3 - iSF_3$ can be aligned nearly in the direction of a primitive $(2,1)$-form, for which S and Z_α get

superheavy masses not far below M_{Pl}, while the gravitino remains to be light. In the second step, one introduces nonperturbative dynamics, e.g. gaugino condensation [7], to stabilize the Kähler moduli T_i. This step fixes T_i at an $N = 1$ supersymmetric AdS vacuum with $m_T \approx m_{3/2} \ln(M_{Pl}/m_{3/2})$ and the vacuum energy density $V_0 \approx -3m_{3/2}^2 M_{Pl}^2$. The last step is to introduce an anti-D3 brane ($\bar{D}3$) providing a positive uplifting potential which would make the total vacuum energy density to be positive but small as desired. $\bar{D}3$ induces also a SUSY breaking vacuum shift which eventually generates the soft SUSY breaking terms of visible fields [8, 9].

The structure of soft terms in KKLT flux compactification has been studied in [8, 9]. It has been noted that such compactification typically leads to $F^T/T \sim m_{3/2}/4\pi^2$ (or even smaller in some special case) and $F^{S,Z} \ll F^T$, implying that the loop-induced anomaly mediation [4] generically provides an important contribution to soft terms. If the visible gauge fields originate from $D3$ branes, the resulting soft terms are dominated by the anomaly mediation whose phenomenology has been extensively studied before [10]. However in KKLT set-up, it is difficult to stabilize the position moduli of $D3$ branes. Also the pure anomaly mediation suffers from the negative slepton mass-square problem. In view of these difficulties, a more attractive possibility is that the visible gauge fields originate from $D7$ branes wrapping a 4-cycle. In such case, the soft terms are induced by a specific mixed modulus-anomaly mediation in which the two mediations give comparable contributions [8, 9].

In fact, the KKLT set-up can be considered a specific example of more general scenario in which the light moduli are stabilized by nonperturbative dynamics yielding an $N = 1$ SUSY AdS vacuum, and this SUSY AdS vacuum is lifted to a SUSY-breaking Minkowski (or de Sitter) vacuum by an appropriate uplifting mechanism which is assumed to be sequestered from the visible sector. Such scenario might be realized also in some class of brane models. Indeed the radion stabilization in 5D orbifold SUGRA based on the boundary and bulk gaugino condensations [11] can provide another example of this scenario which results in the mixed radion-anomaly mediation.

In this paper we wish to examine some phenomenological aspects of this mixed modulus-anomaly mediation, including the pattern of low energy sparticle spectrums and the possibility of electroweak symmetry breaking. The discussion will be made within the framework of 4D effective SUGRA with a SUSY breaking uplifting potential which is sequestered from the visible sector. As we will see, depending upon the anomaly to modulus mediation ratio $\alpha = m_{3/2}/[M_0 \ln(M_{Pl}/m_{3/2})]$ at the GUT scale M_{GUT}, where $M_0 = F^T/(T + T^*)$ denotes the modulus-mediated contribution to the gaugino mass at M_{GUT}, the model can lead to a highly distinctive pattern of superparticle masses at low energy scales. This is essentially due to that the low energy soft parameters in a mixed modulus-anomaly mediation with messenger scale Λ are (approximately) same as those of the pure modulus-mediation with a *mirage messenger scale* $\sim (m_{3/2}/M_{Pl})^{\alpha/2}\Lambda$. The minimal KKLT model predicts $\alpha = 1$, so has a mirage messenger scale close to the intermediate scale $\sqrt{m_{3/2}M_{Pl}}$, while the string, compactification and gauge unification scales are all close to M_{Pl}. If $\alpha = 2$, a striking pattern of low energy superparticle spectrum emerges since the mirage messenger scale is close to TeV: soft masses appear to be unified at TeV although the gauge couplings are unified

at $10^{16} - 10^{17}$ GeV! Although no string theory realization is found yet, $\alpha = 2$ can be naturally obtained by an uplifting mechanism to yield an uplifting potential $V_{\rm lift} \propto 1/(T+T^*)$ [8, 9]. ($\bar{D}3$ in the KKLT set-up gives $V_{\rm lift} \propto 1/(T+T^*)^2$.) Alternatively, one might be able to obtain a somewhat wide range of α (including $\alpha = 2$) by tuning the form of the non-perturbative superpotential [9]. In the next section, we discuss some features of the soft terms in mixed modulus-anomaly mediation, which is largely based on the results of [8, 9]. In sec. 3, we examine the resulting low energy soft parameters and present the results of our phenomenological analysis. Sec. 4 is the conclusion.

2. Soft terms in mixed modulus-anomaly mediation

To make a motivation for our study, let us start with a brief discussion of soft terms in KKLT flux compactification following [8, 9]. In KKLT models on CY orientifold, the dilaton S and complex structure moduli Z_α generically get superheavy masses of the order of compactification scale by the 3-form NS and RR fluxes. This step of stabilizing S and Z_α is assumed to preserve the (approximate) $N=1$ SUSY, so the gravitino remains to be light with $m_{3/2} \ll m_{S,Z}$. To fix the Kähler moduli T_i, one introduces a superpotential of the form $W_{np} = Ae^{-aT_i}$ induced by gaugino condensations on $D7$ branes. Since $m_{Z,S} \gg m_{3/2}$, the stabilization of T_i and also the low energy SUSY breaking can be described by an effective SUGRA obtained after integrating out S and Z_α. For simplicity, here we consider only the case with single Kähler modulus T as the generalization to multi Kähler moduli is rather straightforward. Then the effective SUGRA of T and the gauge and matter superfields on $D7/D3$ can be written as

$$\begin{aligned} S_{N=1} = \int d^4x \sqrt{g^C} \Bigg[\int d^4\theta \, CC^* \left(-3\exp(-K_{eff}/3)\right) \\ + \left\{ \int d^2\theta \left(\frac{1}{4} f_a W^{a\alpha} W^a_\alpha + C^3 W_{eff} \right) + \text{h.c.} \right\} \Bigg], \end{aligned} \tag{2.1}$$

where

$$\begin{aligned} K_{eff} &= K_0(T+T^*) + Z_i(T+T^*) Q_i^* Q_i, \\ W_{eff} &= W_0(T) + \frac{1}{6}\lambda_{ijk} Q_i Q_j Q_k. \end{aligned} \tag{2.2}$$

Here $g^C_{\mu\nu}$ is the 4D metric in superconformal frame which is related to the Einstein frame metric $g^E_{\mu\nu}$ as $g^C_{\mu\nu} = (CC^*)^{-1} e^{K_{eff}/3} g^E_{\mu\nu}$, $C = C_0 + F^C\theta^2$ is the chiral compensator superfield of 4D $N=1$ SUGRA, and Q_i are the gauge-charged matter superfields.

The modulus Kähler and superpotential of the minimal KKLT set-up are given by

$$\begin{aligned} K_0 &= -3\ln(T+T^*), \\ W_0 &= w_0 - Ae^{-aT}, \end{aligned} \tag{2.3}$$

where the constant piece w_0 of the superpotential originates from the fluxes. Using the following $U(1)_R$ transformation of the superconformal formulation of the 4D $N=1$ SUGRA:

$$U(1)_R: \quad C \to e^{i\beta_R} C, \quad W_{eff} \to e^{-3i\beta_R} W_{eff}, \tag{2.4}$$

one can make w_0 to be a real positive constant without loss of generality. The Kähler potential (and also the uplifting potential which will be introduced later) of the KKLT model possesses an approximate nonlinear PQ symmetry,

$$U(1)_T: \quad T \to T + i\beta_T, \tag{2.5}$$

with which one can make A to be a real positive constant again without loss of generality. As was noticed before [9, 12], this nonlinear PQ symmetry is crucial for the KKLT set-up to avoid dangerous SUSY CP violation.

The holomorphic Yukawa couplings λ_{ijk} are independent of T, however the matter Kähler metric and holomorhic gauge kinetic functions can have nontrivial T-dependence as

$$\begin{aligned} Z_i &= \frac{1}{(T+T^*)^{n_i}}, \\ f_a &= T^{\,l_a}, \end{aligned} \tag{2.6}$$

where $n_i = 0$ and $l_a = 1$ for the matter and gauge fields living on $D7$, while $n_i = 1$ and $l_a = 0$ for the matter and gauge fields on $D3$ [13]. When the matter fields live on the intersections of $D7$ branes, n_i can have a fractional value, e.g. $n_i = 1/2$ [13, 14].

The modulus superpotential W_0 stabilizes T at

$$\langle aT \rangle \approx \ln(A/w_0) \approx \ln(M_{Pl}/m_{3/2}), \tag{2.7}$$

however the resulting ground state is a SUSY preserving AdS vacuum. To obtain a SUSY-breaking Minkowski (or dS) vacuum, KKLT proposed to add an anti-$D3$ brane ($\bar{D}3$) providing a positive uplifting potential. Such $\bar{D}3$ is stabilized at the tip of a smoothed conifold singularity at which the geometry is highly warped with an exponentially small warp factor $e^{A_{min}} \sim \sqrt{m_{3/2}/M_{PL}}$ [15]. On $\bar{D}3$, $N = 1$ SUSY is broken explicitly or is non-linearly realized [16]. It has been argued [9] that the low energy consequence of $\bar{D}3$ in KKLT set-up can be described by a single spurion operator up to small corrections further suppressed by $e^{A_{min}}$:

$$S_{\text{lift}} = -\int d^4x \sqrt{g^C} \int d^4\theta \; C^2 C^{2*} \theta^2 \bar{\theta}^2 \; \mathcal{P}_{\text{lift}}(T, T^*), \tag{2.8}$$

where

$$\mathcal{P}_{\text{lift}} = D(T+T^*)^{n_P} \tag{2.9}$$

for a positive constant $D = \mathcal{O}(e^{4A_{min}} M_{Pl}^4) = \mathcal{O}(m_{3/2}^2 M_{Pl}^2)$. Including this spurion operator, the low energy effective action of KKLT compactification is given by

$$S_{eff} = S_{N=1} + S_{\text{lift}}. \tag{2.10}$$

From this, one finds that the SUSY breaking order parameters (in the Einstein frame) approximately take the standard $N = 1$ SUGRA form:

$$\begin{aligned} \frac{F^C}{C_0} &= \frac{1}{3}\partial_T K_0 F^T + e^{K_0/2} W_0^*, \\ F^T &= -e^{K_0/2} K_0^{TT^*} \left(D_T W_0\right)^*, \\ m_{3/2} &= e^{K_0/2} W_0, \end{aligned} \tag{2.11}$$

while the modulus potential contains the uplifting term

$$V_0 = e^{K_0}\left(K_0^{TT^*} D_T W_0 (D_T W_0)^* - 3|W_0|^2\right) + V_{\text{lift}}, \tag{2.12}$$

where

$$V_{\text{lift}} = e^{2K_0/3}\mathcal{P}_{\text{lift}}(T,T^*) \equiv \frac{D}{(T+T^*)^{2-n_P}}. \tag{2.13}$$

It is now straightforward to compute the SUSY breaking order parameters F^T and F^C (or equivalently $m_{3/2}$) by minimizing the above modulus potential under the fine tuning for $\langle V_0 \rangle = 0$. One finds [8, 9]

$$\begin{aligned} \frac{F^C}{C_0} &\approx m_{3/2} \approx \frac{w_0}{M_{Pl}^2 (T+T^*)^{3/2}}, \\ \frac{F^T}{(T+T^*)} &\approx \frac{2-n_P}{a(T+T^*)} m_{3/2}. \end{aligned} \tag{2.14}$$

For the uplifting potential originating from $\bar{D}3$, the corresponding $\mathcal{P}$ is a T-independent constant, i.e. $n_P = 0$ [17, 18].

One of the most interesting features of the KKLT flux compactification is that

$$M_0 \equiv \frac{F^T}{(T+T^*)} = \mathcal{O}\left(\frac{m_{3/2}}{4\pi^2}\right) \tag{2.15}$$

for $m_{3/2}$ near the TeV scale, which suggests that the loop-induced anomaly mediation and the tree-level modulus mediation can be comparable to each other. This is essentially due to the relation $aT \approx \ln(M_{Pl}/m_{3/2}) = \mathcal{O}(4\pi^2)$. As we will see, phenomenological consequences of the mixed modulus-anomaly mediation are somewhat sensitive to the ratio between the anomaly and modulus mediations which we will parameterize by

$$\alpha \equiv \frac{m_{3/2}}{M_0 \ln(M_{Pl}/m_{3/2})} \approx \frac{2}{a(T+T^*)} \frac{F^C}{C_0} \frac{(T+T^*)}{F^T}. \tag{2.16}$$

The minimal KKLT set-up described by the modulus Kähler and superpotential (2.3) and the uplifting potential (2.13) with $n_P = 0$ predicts

$$\alpha\big|_{KKLT} = 1 + \mathcal{O}\left(\frac{1}{\ln(M_{Pl}/m_{3/2})}\right). \tag{2.17}$$

However one might be able to generalize the model to obtain a different value of α. It has been noticed [9] that a model with racetrack superpotential $W_0 = -A_1 e^{-a_1 T} + A_2 e^{-a_2 T}$ [19] can give a stable Minkowski vacuum with a light gravitino and $\alpha = \mathcal{O}(4\pi^2)$ when $(a_2 - a_1)/(a_2 + a_1) = \mathcal{O}(1/4\pi^2)$. Motivated by this observation, one can consider a more general class of effective SUGRA described by

$$\begin{aligned} K_0 &= -n_0 \ln(T+T^*), \\ W_0 &= w_0 - A_1 e^{-a_1 T} + A_2 e^{-a_2 T} \quad (a_1 \leq a_2), \\ \mathcal{P}_{\text{lift}} &= D\,(T+T^*)^{n_P}, \end{aligned} \tag{2.18}$$

for which

$$V_{\rm lift} = e^{2K_0/3}\mathcal{P}_{\rm lift} = \frac{D}{(T+T^*)^{\frac{2}{3}n_0-n_P}}. \tag{2.19}$$

For the parameter regions which give a stable Minkowski vacuum with light gravitino, we have examined the value of α predicted by this model, and found

$$\alpha \,=\, \frac{\xi}{1-3n_P/2n_0}\,, \tag{2.20}$$

where ξ is close to 1 in most of the parameter spaces, but can be significantly bigger than 1 for $|w_0/A_1e^{-a_1T}| \ll 1$ and $(a_2-a_1)/(a_2+a_1) = \mathcal{O}(1/4\pi^2)$ as anticipated in [9]. If the uplifting spurion operator $\mathcal{P}_{\rm lift}$ is a T-independent constant as in the KKLT case, i.e. $n_P = 0$, we have $\alpha = \xi$. Although no concrete realization is found yet, string theory might be able to provide other forms of $\mathcal{P}_{\rm lift}$, e.g. $n_P = 1$ or $n_P = -1$ which gives $\alpha = 2\xi$ or $\alpha = 2\xi/3$ for $n_0 = 3$. In this paper, we simply take this possibility without questioning the origin of $\mathcal{P}_{\rm lift}$, and treat α as a free parameter while focusing on $\alpha = \mathcal{O}(1)$ for which the resulting phenomenology is most interesting. Note that $\alpha \ll 1$ corresponds to the limit of pure modulus-mediation, while $\alpha \gg 1$ corresponds to the pure anomaly mediation.

Let us now consider the soft SUSY breaking terms of canonically normalized visible fields for generic value of α:

$$\mathcal{L}_{soft} \,=\, -\frac{1}{2}M_a\lambda^a\lambda^a - m_i^2|\tilde{Q}_i|^2 - \frac{1}{6}A_{ijk}y_{ijk}\tilde{Q}_i\tilde{Q}_j\tilde{Q}_k + {\rm h.c.}, \tag{2.21}$$

where λ^a are gauginos, $\tilde{Q}_i$ are sfermions, and y_{ijk} denote the canonically normalized Yukawa couplings:

$$y_{ijk} = \frac{\lambda_{ijk}}{\sqrt{e^{-K_0}Z_iZ_jZ_k}}. \tag{2.22}$$

For $\alpha = \mathcal{O}(1)$, the loop-induced anomaly-mediation [4] becomes comparable to the modulus mediation, thus should be included in the soft terms at scales below the compactification (unification) scale. This results in a mixed modulus-anomaly mediation, yielding the following form of soft parameters at energy scale just below the unification scale [8, 9]:

$$\begin{aligned}
M_a \,&=\, F^T\partial_T\ln\left({\rm Re}(f_a)\right) + \frac{b_ag_a^2}{8\pi^2}\frac{F^C}{C_0} \\
&= l_aM_0 + \frac{b_a}{8\pi^2}g_{GUT}^2m_{3/2}, \\
A_{ijk} \,&=\, -F^T\partial_T\ln\left(\frac{\lambda_{ijk}}{e^{-K_0}Z_iZ_jZ_k}\right) - \frac{1}{16\pi^2}(\gamma_i+\gamma_j+\gamma_k)\frac{F^C}{C_0} \\
&= a_{ijk}M_0 - \frac{1}{16\pi^2}(\gamma_i+\gamma_j+\gamma_k)m_{3/2}, \\
m_i^2 \,&=\, \frac{2}{3}V_0 - F^TF^{T*}\partial_T\partial_{T^*}\ln\left(e^{-K_0/3}Z_i\right) - \frac{1}{32\pi^2}\frac{d\gamma_i}{d\ln\mu}\left|\frac{F^C}{C_0}\right|^2
\end{aligned}$$

$$
\begin{aligned}
&+\frac{1}{16\pi^2}\left\{(\partial_T\gamma_i)F^T\left(\frac{F^C}{C_0}\right)^* + \text{h.c.}\right\} \\
&= c_i|M_0|^2 - \frac{1}{32\pi^2}\frac{d\gamma_i}{d\ln\mu}|m_{3/2}|^2 \\
&+\frac{1}{8\pi^2}\left\{\sum_{jk} a_{ijk}\left|\frac{y_{ijk}}{2}\right|^2 - \sum_A l_A g_A^2 C_A(Q_i)\right\}\left(M_0 m_{3/2}^* + M_0^* m_{3/2}\right), \qquad (2.23)
\end{aligned}
$$

where

$$
a_{ijk} = 3 - n_i - n_j - n_k, \quad c_i = 1 - n_i
$$

for the matter Kähler metric $Z_i = 1/(T+T^*)^{n_i}$,

$$
C_A(Q_i)\mathbf{1} = \sum_{a\in A} T_a^2(Q_i)
$$

for the A-th gauge group, and $M_0 = F^T/(T+T^*)$. Here b_a and γ_i are the one-loop beta function coefficients and the anomalous dimension of Q_i, respectively, defined by $\frac{dg_a}{d\ln\mu} = \frac{b_a}{8\pi^2}g_a^3$ and $\frac{d\ln Z_i}{d\ln\mu} = \frac{1}{8\pi^2}\gamma_i$:

$$
\begin{aligned}
b_a &= -\frac{3}{2}\text{tr}\left(T_a^2(\text{Adj})\right) + \frac{1}{2}\sum_i \text{tr}\left(T_a^2(Q_i)\right), \\
\gamma_i &= 2\sum_A g_A^2 C_A(Q_i) - \frac{1}{2}\sum_{jk}|y_{ijk}|^2, \\
\partial_T\gamma_i &= -\frac{1}{2}\sum_{jk}|y_{ijk}|^2\partial_T\ln\left(\frac{\lambda_{ijk}}{e^{-K_0}Z_iZ_jZ_k}\right) - 2\sum_A g_A^2 C_A(Q_i)\partial_T\ln\left(\text{Re}(f_A)\right).
\end{aligned}
$$

Here we have ignored the off-diagonal terms of $\omega_{ij} = \sum_{kl} y_{ikl}y_{jkl}^*$.

If the visible gauge fields originate from $D3$ branes, so $l_a = 0$ and $n_i = 1$, the resulting soft terms correspond to the pure anomaly mediation whose phenomenology has been extensively studied before [10]. However in KKLT set-up, it is difficult to stabilize the position moduli of those $D3$ branes. Also the pure anomaly mediation suffers from the negative slepton mass-square problem. In view of these difficulties, a more attractive possibility is that the visible gauge fields originate from $D7$ branes, for which $l_a = 1$ but still n_i can be either 0 or 1/2 or even 1, depending on the origin of Q_i. In the following, we will set $l_a = 1$, and then the generic mixed modulus-anomaly mediation is parameterized by

$$
M_0,\ a_{ijk},\ c_i,\ \alpha = m_{3/2}/[M_0\ln(M_{Pl}/m_{3/2})],
$$

where the first three parameters are determined by the modulus-dependence of the matter Kähler metric, while α is determined by the mechanism of modulus stabilization and the subsequent uplifting.

For the minimal KKLT set-up defined by (2.3), (2.6) and (2.9), the invariance of the matter Kähler metric (2.6) under the nonlinear PQ symmetry (2.5) assures that a_{ijk}

are real. We already noticed that the $U(1)_R$ transformation (2.4) and the nonlinear PQ transformation (2.5) can be used to make the two parameters w_0 and A in the modulus superpotential (2.3) to be real without loss of generality. In such field basis, the resulting F^T and F^C are real, thus the gaugino masses and A-parameters in the minimal KKLT set-up do not contain any dangerous CP-violating phase [9, 12].

For a later discussion of the electroweak symmetry breaking, let us discuss the Higgs mass parameters μ and B in the mixed modulus-anomaly mediation scenario. (Here we are using the same notation μ for both the Higgsino mass parameter and the renormalization point.) One possible source of μ and B is the Higgs bilinear terms in Kähler and superpotential, which would take the following form [20, 21]

$$\begin{aligned}\Delta K_{eff} &= \frac{H_1 H_1^*}{(T+T^*)^{n_{H_1}}} + \frac{H_2 H_2^*}{(T+T^*)^{n_{H_2}}} + \left(\frac{\kappa H_1 H_2}{(T+T^*)^h} + \text{h.c.}\right), \\ \Delta W_{eff} &= \tilde{A} e^{-aT} H_1 H_2, \end{aligned} \tag{2.24}$$

where ΔW_{eff} is induced by the non-perturbative dynamics yielding Ae^{-aT} in the modulus superpotential (2.3). The resulting μ for the canonically normalized Higgs doublets at the unification scale is given by

$$\mu = \mu_W + \mu_K \,, \tag{2.25}$$

where

$$\mu_W = \frac{\tilde{A} e^{-aT}}{(T+T^*)^{l_W}}, \quad \mu_K = \frac{\kappa}{(T+T^*)^{l_K}} \left(\frac{F^C}{C_0} + (1-h)\frac{F^T}{(T+T^*)}\right)^* \tag{2.26}$$

for $l_W = (3 - n_{H_1} - n_{H_2})/2$ and $l_K = (2h - n_{H_1} - n_{H_2})/2$. Soft masses in the mixed modulus-anomaly mediation are of the order of $M_0 = F^T/(T+T^*)$, thus it is desirable that μ is $\mathcal{O}(M_0)$ also. Although $F^C/C_0 \approx m_{3/2} = \mathcal{O}(4\pi^2 M_0)$ in the mixed modulus-anomaly mediation, it is not difficult to obtain $\mu = \mathcal{O}(M_0)$. In view of that $\tilde{A} \approx A/M_{Pl}^2$ and $Ae^{-aT} \approx m_{3/2} M_{Pl}^2/aT$ for the modulus superpotential (2.3), μ_W is naturally of the order of M_0. As for μ_K, one might assume that the $H_1 H_2$-term in K_{eff} is induced by a loop correction, and thus $\mu_K = \mathcal{O}(m_{3/2}/4\pi^2) = \mathcal{O}(M_0)$.

In fact, the real problem is to get the desired size of B. For B originating from (2.24), we find

$$\begin{aligned} B\mu &= -\left[m_{3/2} - a(T+T^*)M_0 + \mathcal{O}(M_0)\right]\mu_W + \left[m_{3/2} + \mathcal{O}(M_0)\right]\mu_K \\ &= m_{3/2}\left[\,\mu_K + \left(\frac{2}{\alpha} - 1\right)\mu_W\,\right] + \mathcal{O}(M_0^2) \end{aligned} \tag{2.27}$$

where we have used $\alpha \approx 2m_{3/2}/a(T+T^*)M_0$. This result shows that B from (2.24) is generically of the order of $m_{3/2}$, thus would be too large to achieve the correct electroweak symmetry breaking. Unless $\alpha \approx 1$, one can still obtain the desired $B = \mathcal{O}(M_0)$ under the fine-tuning:

$$\mu_K + \left(\frac{2}{\alpha} - 1\right)\mu_W = \mathcal{O}\left(\frac{\mu_K + \mu_W}{4\pi^2}\right). \tag{2.28}$$

which we will assume in the discussion of electroweak symmetry breaking in the next section. We note that for the models yielding $\alpha \approx 2$, for instance a model with $n_0 = 3, n_P = 1, A_2 = 0$ in (2.18) which gives $\alpha = 2 + \mathcal{O}(1/4\pi^2)$, the above condition is automatically satisfied when $\mu_K = 0$, i.e. when the μ term originates entirely from the nonperturbative term in W_{eff}. For the case with $\alpha \approx 1$, which would be the most interesting case as it is predicted by the minimal KKLT set-up, the fine-tuning (2.28) is not allowed. One might then consider

$$\begin{aligned}\Delta K_{eff} &= \frac{H_1 H_1^*}{(T+T^*)^{n_{H_1}}} + \frac{H_2 H_2^*}{(T+T^*)^{n_{H_2}}} + \left(\frac{\kappa H_1 H_2}{(T+T^*)^h} + \text{h.c.}\right), \\ \Delta W_{eff} &= \tilde{\mu} H_1 H_2,\end{aligned} \tag{2.29}$$

where $\tilde{\mu}$ is a T-independent constant which is adjusted to give $\mu_W = \tilde{\mu}/(T+T^*)^{l_W} = \mathcal{O}(M_0)$. In this case, one easily finds $\mu = \mu_K + \mu_W$ and $B\mu = m_{3/2}(\mu_K - \mu_W) + \mathcal{O}(M_0^2)$, thus $B = \mathcal{O}(M_0)$ can be obtained through the fine-tuning: $\mu_K - \mu_W = \mathcal{O}(\mu/4\pi^2)$.

Generating the μ and B-terms through (2.24) or (2.29) has another unattractive feature in addition to the involved fine-tuning to get $B = \mathcal{O}(M_0)$. Those models for μ and B involve a CP-violating phase $\text{Arg}(\tilde{A}\kappa^*)$ or $\text{Arg}(\tilde{\mu}\kappa^*)$ which eventually generates a nonzero $\text{Arg}(B)$. To satisfy the constraints from the hadron or electron electric dipole moments, one then needs to tune this CP phase to be smaller than 10^{-2}.

The difficulty to obtain $B = \mathcal{O}(M_0)$ is a generic problem of models which predict $m_{3/2} \gg M_0$. A simple way to avoid this difficulty is to assume that the Higgs μ-term originates from a trilinear Yukawa term involving a singlet N [22]:

$$\Delta W_{eff} = \lambda_1 N H_1 H_2 + \frac{\lambda_2}{3} N^3. \tag{2.30}$$

In this case, the (effective) μ and B are given by

$$\begin{aligned}\mu &= e^{K_0/2} \lambda_1 \langle N \rangle, \\ B &= A_{N H_1 H_2} + e^{K_0/2} \lambda_2^* \langle N \rangle / Z_N,\end{aligned} \tag{2.31}$$

where Z_N is the Kähler metric of N, thus they have the desired size of $\mathcal{O}(M_0)$. In addition to giving the correct size of μ and B without a fine-tuning of parameters, this model has another virtue: it avoids naturally all the potentially dangerous CP phases in the soft parameters. Using the field redefinitions

$$U(1)_H : H_1 H_2 \to e^{i\beta_H} H_1 H_2, \quad U(1)_N : N \to e^{i\beta_N} N, \tag{2.32}$$

one can make λ_1 and λ_2 to be real, for which $\langle N \rangle$ is real also. We already noticed that the field redefinitions of (2.4) and (2.5) can be used to assure that the gaugino masses and A-parameters in the minimal KKLT set-up are all real. When the Higgs μ and B parameters originate from the superpotential (2.30), the resulting B is automatically real, thus the model is completely free from the potentially dangerous SUSY CP violation.

Before closing this section, we note that the mixed modulus-anomaly mediation can arise also from 5D brane models on S^1/Z_2 stabilizing the radion by nonperturbative effects

[11]. To see this, let us first note that gauge fields propagating in 5D bulk have $f_a = T$ for the radion superfield $T = R + iB_5$ where R is the orbifold radius and B_5 is the fifth component of the 5D graviphoton [23]. On the other hand, for gauge fields confined in the boundaries of S^1/Z_2, the corresponding f_a are T-independent constants. The radion Kähler potential is given by $K_0 = -3\ln(T+T^*)$, and a superpotential of the form $W_0 = w_0 - \sum_i A_i e^{-a_i T}$ can be generated by gaugino condensations. The constant piece w_0 would be induced by a gaugino condensation at the boundary, while $\sum_i A_i e^{-a_i T}$ arises from bulk gaugino condensations. In the simplest situation, the matter Kähler metric takes the form of (2.6), $n_i = 1$ for boundary matter fields and $n_i = 0$ for bulk matter fields [23]. In fact, if a bulk matter field Q_i have nonzero 5D mass M_i, its zero mode has a Kähler metric given by $Z_i = (1 - e^{M_i(T+T^*)})/M_i(T+T^*)$ [24]. In such case, the resulting Yukawa couplings $y_{ijk} = \lambda_{ijk}/\sqrt{e^{-K_0} Z_i Z_j Z_k}$ can have a hierarchical structure in a natural manner. Again the radion superpotential W_0 from gaugino condensations can stabilize T at a SUSY AdS vacuum. Adding a SUSY-breaking anti-brane at the boundary will uplift this AdS vacuum to a Minkowski (or dS) vacuum. The corresponding spurion operator $\mathcal{P}_{\text{lift}}$ is a T-independent constant, so $n_P = 0$ [11]. If $\mathcal{P}_{\text{lift}}$ originates dominantly from a T-dependent bulk $U(1)$ FI-term $D_{FI} \propto \partial_T K_0$ [17], one would have $\mathcal{P}_{\text{lift}} = e^{-2K_0/2} V_D \propto g^2 e^{-2K_0/3} (\partial_T K_0)^2 \propto 1/(T+T^*)$ so $\alpha \approx 2/3$. On the other hand, if $\mathcal{P}_{\text{lift}}$ originates dominantly from a T-independent $U(1)_R$ FI term [25], $\mathcal{P}_{\text{lift}} \propto g^2 e^{-2K_0/3} \propto (T+T^*)$ and then $\alpha \approx 2$. However, as was pointed out in [9], $U(1)$ FI-dominated uplifting is difficult to be achieved since the FI D-term always vanishes for a SUSY AdS solution of the F-term potential.

3. Low energy phenomenology

In this section, we examine some phenomenological consequences of the mixed modulus-anomaly mediation given by (2.23) with $l_a = 1$ and α in the range of $\mathcal{O}(1)$. For concrete analysis, we will use the standard value of the unification scale $M_{GUT} \sim 2 \times 10^{16}$ GeV.

Mixed modulus-anomaly mediation can give a low energy sparticle spectrum which is quite different from other scenarios of SUSY breaking. This is mainly due to the particular correlation between the anomaly mediation and the RG evolution of soft parameters. To see this, let us consider the low energy gaugino masses. At a scale $M_{GUT}^{(-)}$ just below M_{GUT}, the gaugino masses are given by

$$M_a(M_{GUT}^{(-)}) = M_0 \left(1 + \frac{\ln(M_{Pl}/m_{3/2})}{8\pi^2} \, \alpha b_a g_{GUT}^2 \right), \tag{3.1}$$

where α is defined in (2.16). The one-loop RG evolution between $M_{GUT}^{(-)}$ and μ yields

$$M_a(\mu) = M_0 \left[1 - \frac{1}{4\pi^2} b_a g_a^2(\mu) \ln \left(\frac{M_{GUT}}{(M_{Pl}/m_{3/2})^{\alpha/2} \mu} \right) \right]. \tag{3.2}$$

This result shows an interesting feature: if $\alpha = \mathcal{O}(1)$ as in KKLT set-up, the anomaly-mediated contribution, i.e. the α-dependent part, cancels significantly the RG evolution of

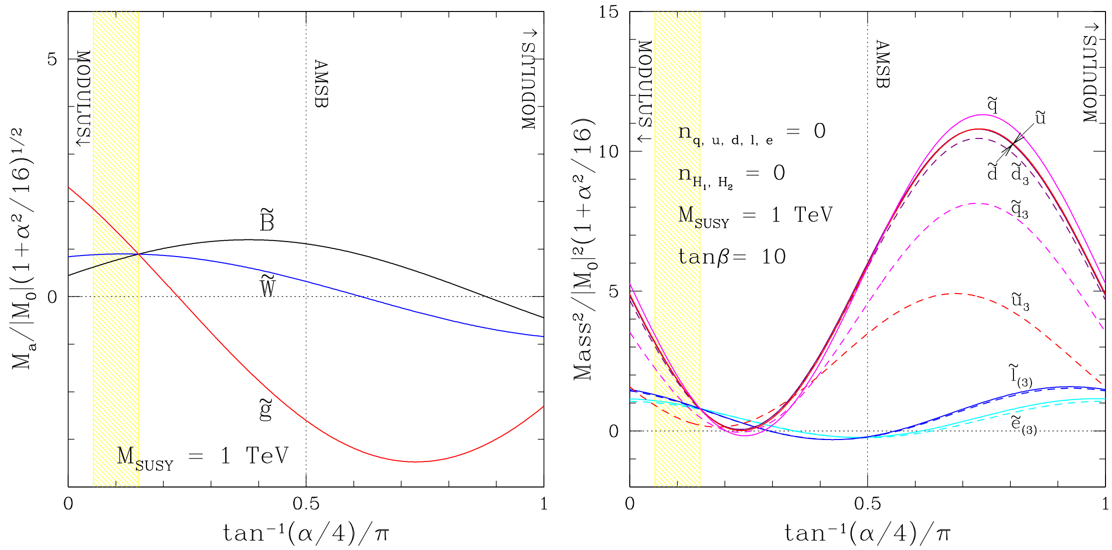

Figure 1: Sparticle mass spectrum (relative ratios) at $M_{SUSY} = 1$ TeV for the entire range of the anomaly to modulus ratio $\alpha = m_{3/2}/[M_0 \ln(M_{Pl}/m_{3/2})]$. Here $M_0 = F^T/(T+T^*)$ and $m_{3/2} = F^C/C_0$. The shaded region indicates the range $2/3 \leq \alpha \leq 2$ and the short-dashed curves denote the 3rd generation squarks/sleptons. Note that the sign convention of the gaugino masses (and A_{ijk}) for $0 \leq \tan(\alpha/4) \leq \pi/2$ is different from the convention for $\pi/2 \leq \tan(\alpha/4) \leq \pi$.

the modulus-mediated gaugino masses. In particular it shows that the low energy gaugino masses in the mixed modulus-anomaly mediation started from the messenger scale M_{GUT} are same as the low energy gaugino masses in the pure modulus-mediation started from a *mirage messenger scale* $\sim (m_{3/2}/M_{Pl})^{\alpha/2} M_{GUT}$. Note that this mirage messenger scale does not correspond to a physical threshold scale. Still the physical gauge coupling unification scale is M_{GUT}, and the Kaluza-Klein and string threshold scales are a little above M_{GUT}. When $\alpha \approx 2$, $M_a(M_{SUSY})$ are approximately same as the pure modulus-mediated gaugino mass without any RG running effect,

$$M_a(M_{SUSY}) \approx M_0 \quad \text{for} \quad \alpha \approx 2. \tag{3.3}$$

In this case, we have unified gaugino masses at TeV, while the corresponding gauge couplings are unified at $M_{GUT} \sim 2 \times 10^{16}$ GeV. The low energy values of A_{ijk} and m_i^2 show a similar feature. In fact, if y_{ijk} are non-vanishing *only* for the combinations $Q_i Q_j Q_k$ satisfying $a_{ijk} = 1$ and $c_i + c_j + c_k = 1$, *or* if the effects of the Yukawa couplings on the renormalization group evolution can be ignored, the resulting A_{ijk} and m_i^2 at low energies are given by

$$\begin{aligned} A_{ijk}(\mu) &= M_0 \left[a_{ijk} + \frac{1}{8\pi^2}(\gamma_i(\mu) + \gamma_j(\mu) + \gamma_k(\mu)) \ln \left(\frac{M_{GUT}}{(M_{Pl}/m_{3/2})^{\alpha/2}\mu} \right) \right], \\ m_i^2(\mu) &= |M_0|^2 \left[c_i + \frac{1}{4\pi^2} \left\{ \gamma_i(\mu) - \frac{1}{2}\frac{d\gamma_i(\mu)}{d\ln\mu} \ln \left(\frac{M_{GUT}}{(M_{Pl}/m_{3/2})^{\alpha/2}\mu} \right) \right\} \right. \end{aligned}$$

$$\times \ln\left(\frac{M_{GUT}}{(M_{Pl}/m_{3/2})^{\alpha/2}\mu}\right) - \frac{1}{8\pi^2}Y_i\Big(\sum_j c_jY_j\Big)g_Y^2(\mu)\ln\left(\frac{M_{GUT}}{\mu}\right)\Bigg], \quad (3.4)$$

where $\gamma_i(\mu)$ and $g_Y^2(\mu)$ denote the running anomalous dimensions and the running $U(1)_Y$ gauge coupling at μ, respectively, and Y_i is the $U(1)_Y$ hypercharge of Q_i. Here again we ignored the off-diagonal parts of $\omega_{ij} = \sum_{kl} y_{ikl}y^*_{jkl}$, and the last part of $m_i^2(\mu)$ depending on $\sum_i c_iY_i$ arises as a consequence of the $\mathrm{Tr}(Ym^2)$-term in the RG equation of m_i^2. In generic situation, the above results will be modified by the Yukawa couplings for $a_{ijk} \neq 1$ or $c_i + c_j + c_k \neq 1$. However for the first and second generations of quarks and leptons, the modification will be negligible since the involved Yukawa couplings are small enough. Again the anomaly mediated contributions at M_{GUT} leads to a mirage messenger scale $(m_{3/2}/M_{Pl})^{\alpha/2}M_{GUT}$. Also if $\alpha \approx 2$, the TeV scale sfermion masses are approximately same as the pure modulus-mediated sfermion masses without any RG running effect:

$$m_i^2(M_{SUSY}) \approx c_i\,|M_0|^2 \quad \text{for} \quad \alpha \approx 2 \qquad (3.5)$$

up to ignoring the corrections due to y_{ijk} for $a_{ijk} \neq 1$ or $c_i + c_j + c_k \neq 1$.

The results of (3.2) and (3.4) show that the low energy soft parameters in a mixed modulus-anomaly mediation with messenger scale Λ are approximately same as those of the pure modulus-mediation with a mirage messenger scale $(m_{3/2}/M_{Pl})^{\alpha/2}\Lambda$. We note that this feature does not depend on the detailed form of the modulus-mediation, thus applies to generic form of the mixed modulus-anomaly mediation. As was pointed out in [9], the mixed modulus-anomaly mediation can avoid dangerous SUSY flavor and CP violations in a natural manner. The soft terms preserve the quark and lepton flavors if n_i are flavor-independent, which would arise automatically if the matter fields with common gauge charge originate from the same geometric structure. They also preserve CP since the relative CP phase between F^T and F^C could be rotated away by the shift of the axion-component of T [12]. In Figs. 1−3, we depict the pattern of low energy sparticle masses for an appropriate range of α. In this calculation, we determined the gauge and Yukawa couplings at $M_{SUSY} = 1$ TeV using the 2-loop RG equations of the SM and the top quark mass $m_t^{pole} = 178$ GeV. Above M_{SUSY}, we have used the 1-loop RG equation of the MSSM to arrive at M_{GUT}. The results are not affected significantly by a different choice of M_{SUSY} as long as it is not far from the weak scale. Fig. 1 provides an overall view of the low energy sparticle spectrum for the whole range of $\tan^{-1}(\alpha/4)$, including the pure modulus mediation ($\tan^{-1}(\alpha/4) = 0$) and also the pure anomaly mediation ($\tan^{-1}(\alpha/4) = \pi/2$). Here we assumed that all MSSM matter and Higgs fields originate from $D7$ branes, thus $n_i = 0$ for all MSSM chiral multiplets, and used the top quark Yukawa coupling for $\tan\beta = 10$. Depending upon the value of α, the gluino to Wino mass ratio (and also the squark to slepton mass ratios) can be considerably suppressed (e.g. $\tan^{-1}(\alpha/4) \sim 0.22\pi$) or enhanced (e.g. $\tan^{-1}(\alpha/4) \sim 0.75\pi$), relative to the pure modulus mediation case. Note that the range $0.20\pi \leq \tan^{-1}(\alpha/4) \leq 0.64\pi$ gives either a tachyonic squark or a tachyoinc slepton. For $0.64\pi \leq \tan^{-1}(\alpha/4) \leq 0.77\pi$, the negative slepton mass-square of the pure anomaly mediation can be lifted to a positive value, while keeping the Wino LSP. The relative signs

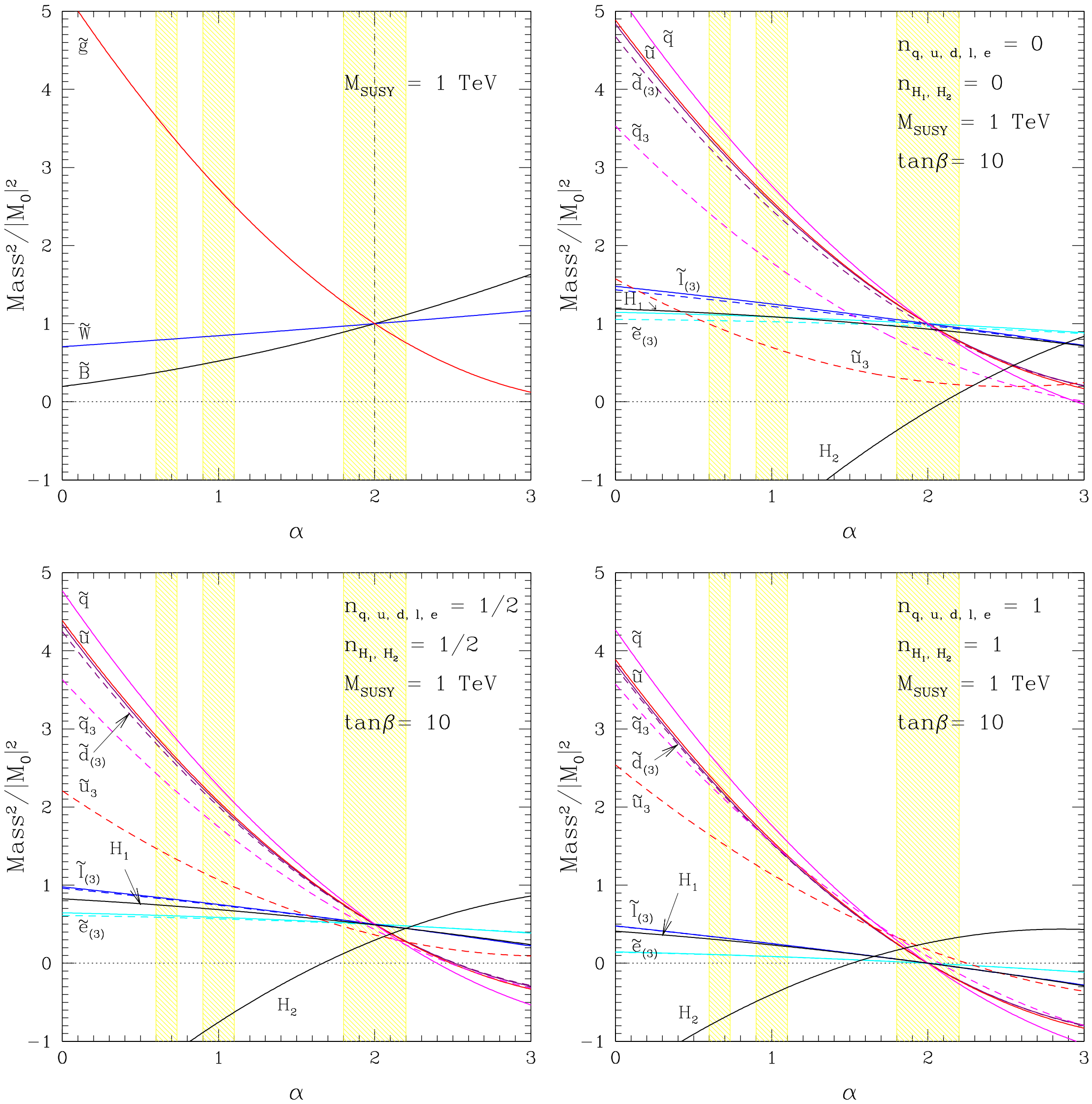

Figure 2: Sparticle mass spectrum at $M_{SUSY} = 1$ TeV for $0 \leq \alpha \leq 3$. The shaded regions correspond to the moduli Kähler and superpotential (2.3) and the uplifting potential (2.13) with $n_P = -1, 0, 1$ ($\alpha = 2/3, 1, 2$), taking into account 10% uncertainty. Again the short-dashed curves denote the 3rd generation sfermions.

of Bino and gluino masses against the Wino mass can be $(+,+)$ or $(-,+)$ since the $(+,-)$ case is excluded due to a tachyonic slepton/squark.

Fig. 2 shows the low energy sparticle masses for $0 \leq \alpha \leq 3$. This range of α contains $\alpha = 1$ predicted by the minimal KKLT set-up and also $\alpha = 2$ which has a mirage messenger scale close to TeV. As for the squark/slepton spectrums, we considered three different cases distinguished by the universal values of n_i: $0, 1/2$ or 1. For $\alpha = 1$ which is predicted by the minimal KKLT set up, we summarized the resulting sparticle spectrums at M_{SUSY} in

Table. 1. Note that the gluino to Wino/Bino mass ratios $M_{\tilde{g}}/M_{\tilde{W},\tilde{B}}$ and also the squark to slepton mass ratios $m_{\tilde{q}}/m_{\tilde{l}}$ for the minimal KKLT value are significantly smaller than the values predicted by the general mSUGRA scenario or the pure anomaly mediation scenario. For $\alpha = 2$, the low energy superparticle masses are approximately unified (up to corrections due to the top quark Yukawa coupling) since it gives a mirage messenger scale close to 1 TeV. Note that for a fixed value of α, the sparticle mass ratios discussed here are insensitive to the overall size of the SUSY breaking order parameter F^T/T as well as of the details of the electroweak symmetry breaking which will be discussed below.

	$\alpha = 1$	$(\alpha = 1)/(\alpha = 0)$
$\tilde{B}$	0.79	1.48
$\tilde{g}$	1.79	0.66

	$\tan\beta = 10$					
	$n_i = 0$		$n_i = 1/2$		$n_i = 1$	
	$\alpha = 1$	$(\alpha = 1)/(\alpha = 0)$	$\alpha = 1$	$(\alpha = 1)/(\alpha = 0)$	$\alpha = 1$	$(\alpha = 1)/(\alpha = 0)$
$\tilde{e}$	1.13	0.89	0.83	0.87	0.32	0.71
$\tilde{e}_3$	1.10	0.90	0.82	0.88	0.32	0.71
$\tilde{\ell}$	1.22	0.84	0.94	0.80	0.55	0.66
$\tilde{\ell}_3$	1.20	0.84	0.94	0.80	0.55	0.66
$\tilde{d}$	1.73	0.66	1.55	0.63	1.35	0.58
$\tilde{d}_3$	1.70	0.66	1.54	0.63	1.35	0.58
$\tilde{u}$	1.74	0.66	1.56	0.63	1.36	0.58
$\tilde{u}_3$	0.91	0.61	1.12	0.64	1.16	0.61
$\tilde{q}$	1.81	0.66	1.64	0.63	1.44	0.59
$\tilde{q}_3$	1.45	0.65	1.44	0.63	1.35	0.60

Table 1: Sparticle spectrum of the minimal KKLT set-up ($\alpha = 1$) at M_{SUSY}. All masses are divided by the Wino mass M_2 at 1 TeV. The ratios to the pure modulus mediation ($\alpha = 0$) are also presented.

Let us now examine the electroweak symmetry breaking in the mixed modulus-anomaly mediation scenario. This issue depends on how the Higgs mass parameters μ and B are generated. An important phenomenological issue which is sensitive to μ is the nature of the lightest supersymmetric particle (LSP). Also if one could have a concrete theoretical scheme to relate the μ and B to $M_0 = F^T/(T + T^*)$, the overall size of M_0 might be constrained by the condition of the correct electroweak symmetry breaking. Here we will restrict the analysis to the simplest (though not the most attractive) scheme to generate μ and B: the minimal supersymmetric standard model (MSSM) with μ and B obtained from (2.24) or (2.29) under appropriate fine-tuning, while leaving the analysis for the next minimal supersymmetric standard model (NMSSM) with a singlet N, i.e. the model of (2.30), for future work. We also limit the analysis to the tree-level Higgs potential for simplicity.

The neutral part of the Higgs potential is given by

$$V = \tilde{m}_1^2|H_1^0|^2 + \tilde{m}_2^2|H_2^0|^2 - \left(B\mu H_1^0 H_2^0 + \text{c.c.}\right) + \frac{1}{8}(g_1^2+g_2^2)\left(|H_1^0|^2 - |H_2^0|^2\right)^2 , \qquad (3.6)$$

where $\tilde{m}_{1,2}^2 = m_{H_{1,2}}^2 + |\mu|^2$. If the Higgs soft masses and μ satisfy the conditions for a symmetry breaking stable vacuum [26]:

$$\tilde{m}_1^2\tilde{m}_2^2 - |B\mu|^2 < 0 , \quad \tilde{m}_1^2 + \tilde{m}_2^2 - 2|B\mu| > 0, \qquad (3.7)$$

we obtain the following relations:

$$\mu^2 = -\frac{M_Z^2}{2} + \frac{m_{H_1}^2 - m_{H_2}^2 \tan^2\beta}{\tan^2\beta - 1} , \quad |B\mu| = \frac{\tan\beta}{1+\tan^2\beta}(m_{H_1}^2 + m_{H_2}^2 + 2\mu^2), \quad (3.8)$$

which allow us to determine μ/M_0 and B/M_0 in terms of M_Z/M_0, $\tan\beta$ and m_{H_i}/M_0. In Fig.3, we plot the resulting μ/M_0, B/M_0 and the LSP mass for various choices of n_i of the squarks, sleptons and Higgs fields which are assumed to have the Kähler metric $Z_i = (T+T^*)^{-n_i}$. As is shown, the qualitative behavior of μ and B is common to all models, while the precise position of the curves and the nature of LSP are sensitive to the choice of n_i. Thin sold curve for μ indicates the value of μ in the limit $M_Z/M_0 \to 0$, while the dashed curve represents μ for $M_Z/M_0 = 0.3$. Positiveness of M_Z^2 in (3.8) ensures that there is no symmetry breaking solution for $|\mu|$ above the thin solid curve. Because the slepton masses are typically $\mathcal{O}(M_0)$, we are required to choose M_0 significantly bigger than M_Z. Note that μ/M_0 is almost independent of M_Z/M_0 for $M_0 \geq 3.3M_Z$, which amounts to the well known fine-tuning of μ required for the correct electroweak symmetry breaking in MSSM [29]. Starting from the pure modulus mediation ($\alpha = 0$), increasing the anomaly mediated contribution eventually erases the y_t-induced radiative correction to $m_{H_2}^2$ and finally restores the electroweak gauge symmetry, which corresponds to the value of α for which $\mu = 0$. The value of $|B|$ blows up at this value of α unless $|B|$ becomes zero before arriving this value of α, which would make the Higgs potential unbounded from below. These features can be understood by noting that the mirage messenger scale is given by $(m_{3/2}/M_{Pl})^{\alpha/2}M_{GUT}$, thus the strength of radiative electroweak symmetry breaking [26, 27] becomes weaker when α increases from 0 to a positive value of $\mathcal{O}(1)$.

When α increases from 0 to a positive value, the lightest neutralino is changed from the Bino-like to the Higgsino-like around the point where μ crosses the Bino mass. Typically, the minimal KKLT value ($\alpha = 1$) corresponds to the Bino or Bino–Higgsino mixing region while $\alpha = 2$ corresponds to an almost pure Higgsino, which will lead to a considerably different consequence in the dark matter scenario relative to the general mSUGRA case if the mixing is sufficiently large ($\alpha = 1$) or Higgsino is not too light ($\alpha = 2$) [28]. The model-dependence of the value of μ originates mainly from the top Yukawa coupling. On the other hand, μ is rather insensitive to $\tan\beta$ except for $\tan\beta \simeq 1$ which is disfavored by the Higgs boson search in LEPII [30]. This is because most of the contribution to μ^2 in (3.8) comes from $m_{H_2}^2$ as the contribution from $m_{H_1}^2$ is strongly suppressed by $\tan^{-2}\beta$. The green curves in Fig.3 indicate the LSP mass in the unit of M_0. The nature of LSP is somewhat model-dependent, and can be neutralino, stau or stop, depending upon the choice of n_i, α

and $\tan\beta$. If the LSP is stable, this feature can provide a strong constraint on the model. In general, a heavier H_2 lowers the stop mass through the radiative correction involving the top Yukawa coupling, while the choice $n_i = 1$ or a large $\tan\beta$ makes the slepton lighter.

As we have noted, if $y_{ijk} \neq 0$ only for $a_{ijk} = 1$ and $c_i + c_j + c_k = 1$ which would be obtained when $n_i + n_j + n_k = 2$, and also if $\alpha \approx 2$, the RG evolution of modulus mediation is almost canceled by the anomaly mediation, which results in the mirage messenger scale close to TeV. This set-up may provide a new insight for the little hierarchy problem in supersymmetric standard model. In Fig.4, we choose $n_i = 1/2$ for matter fields and $n_i = 1$ for the Higgs fields, satisfying the condition that $n_i + n_j + n_k = 2$ for nonzero Yukawa couplings. The modulus-mediated soft mass-squares of the Higgs bosons at M_{GUT} are vanishing up to small threshold corrections of $\mathcal{O}(M_0^2/8\pi^2)$, and thus $m_{H_i}^2(M_{SUSY})/M_0^2 = \mathcal{O}(1/8\pi^2)$ for $\alpha = 2$. One the other hand, the modulus-mediated squark/slepton mass squares and the modulus-mediated gaugino mass at M_{SUSY} are $M_0^2/2$ and M_0, respectively, thus $M_a(M_{SUSY}) \approx M_0$ and $m_{\tilde{q},\tilde{l}}^2(M_{SUSY}) \approx M_0^2/2$ for $\alpha = 2$. This might enable us to generate a little hierarchy between the weak scale and the sparticle mass in a natural manner: $m_H^2/M_0^2 = \mathcal{O}(1/8\pi^2)$. Of course, we then need a mechanism to generate μ and B smaller than M_0 by one order of magnitude. In general, keeping the hierarchy of $\mathcal{O}(10^{-2})$ between the Higgs mass-squares and the gaugino/squark mass-squares is highly non-trivial because the latter enters in the former through the radiative corrections involving the strong coupling constants [31]. However, in the mixed modulus-anomaly mediation with $\alpha \approx 2$, those radiative corrections are automatically canceled by the anomaly-mediated contributions. Note that although no string theory realization is found yet, $\alpha = 2$ can be naturally obtained by an effective theory with uplifting potential $V_{\rm lift} \propto 1/(T + T^*)$.

4. Conclusion

In this paper, we have examined some phenomenological consequences of the mixed modulus-anomaly mediation scenario for supersymmetry breaking in which the modulus mediation and the anomaly mediation give comparable contributions to soft parameters at the messenger scale $\sim M_{GUT}$. Such mediation scheme can arise naturally in compactified string and brane models which stabilize the light moduli by nonperturbative effects at a supersymmetric AdS vacuum and then break SUSY by a sequestered uplifting potential. A concrete example of such scenario has been proposed recently by KKLT [6] and the pattern of resulting soft terms are analyzed in [8, 9]. The scheme may also offer an interesting cosmological scenario which produces the correct amount of the neutralino dark matter while avoiding the cosmological moduli/gravitino problem [32]. Here we considered a more general set-up based on 4D effective SUGRA with SUSY breaking uplifting potential, and noted that the scheme can result in a highly distinctive superparticle spectrum at low energy scales. This can be easily understood by noting that the low energy soft parameters in a mixed modulus-anomaly mediation with messenger scale Λ are (approximately) same as those of the pure modulus-mediation with a *mirage messenger scale* $\sim (m_{3/2}/M_{Pl})^{\alpha/2}\Lambda$ where $\alpha = m_{3/2}/[M_0 \ln(M_{Pl}/m_{3/2})]$ for M_0 denoting the modulus-mediated contribution to the gaugino mass at M_{GUT}. The minimal KKLT model predicts $\alpha = 1$, thus has a mirage mes-

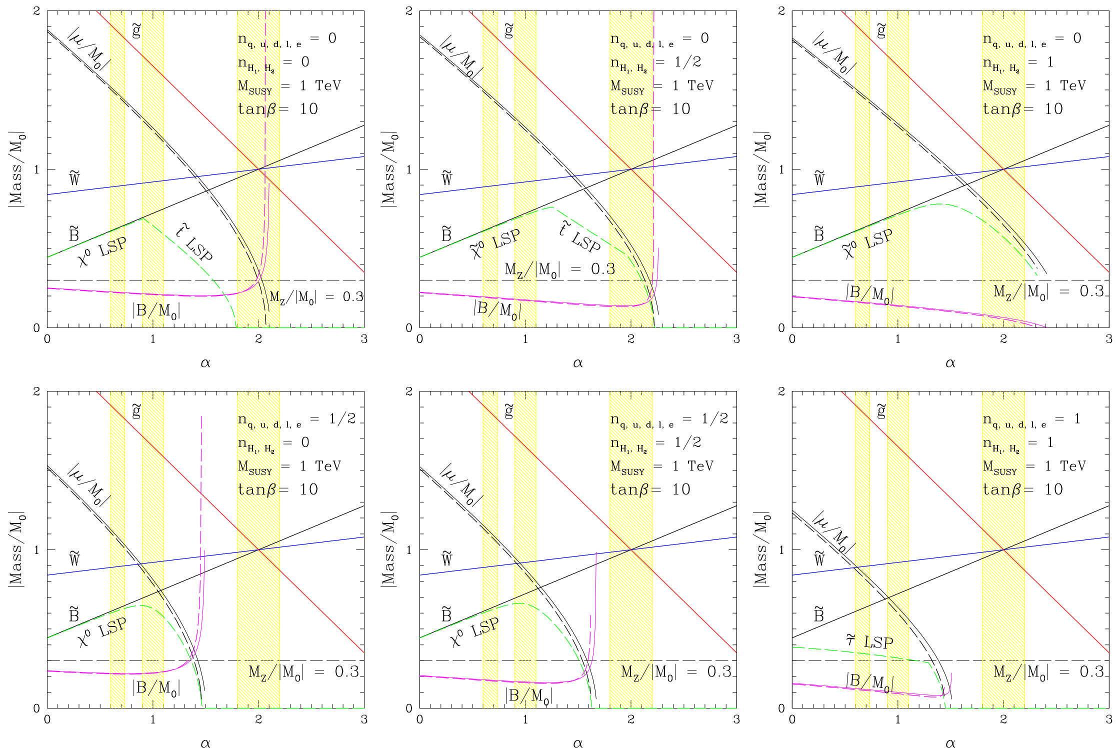

Figure 3: The behavior of the Higgsino mass parameter μ. The shaded region is same as in Fig.2. The dashed (thin–solid) curve corresponds to $M_Z = 0.3M_0$ ($M_0/M_Z \to \infty$).

senger scale close to the intermediate scale $\sqrt{m_{3/2}M_{Pl}}$, while the string, compactification and gauge unification scales are all close to M_{Pl}. The most dramatic situation is $\alpha = 2$ for which soft masses appear to be unified at TeV although the gauge couplings are unified at $10^{16} - 10^{17}$ GeV. Although no string theory realization is found yet, $\alpha = 2$ can be naturally obtained by an uplifting mechanism to yield an uplifting potential $V_{\rm lift} \propto 1/(T+T^*)$ [8, 9]. Alternatively, one might be able to obtain such a value of α by tuning the form of the non-perturbative superpotential [9]. All the results of our phenomenological analysis are summarized in Figs. 1–4 and Table 1.

Acknowledgments

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References

[1] H. P. Nilles, Phys. Rept. **110**, 1 (1984); H. E. Haber and G. L. Kane, Phys. Rept. **117**, 75 (1985).

[2] D. J. H. Chung, L. L. Everett, G. L. Kane, S. F. King, J. Lykken and Lian-Tao Wang , Phys. Rept. **407**, 1 (2005) [arXiv:hep-ph/0312378].

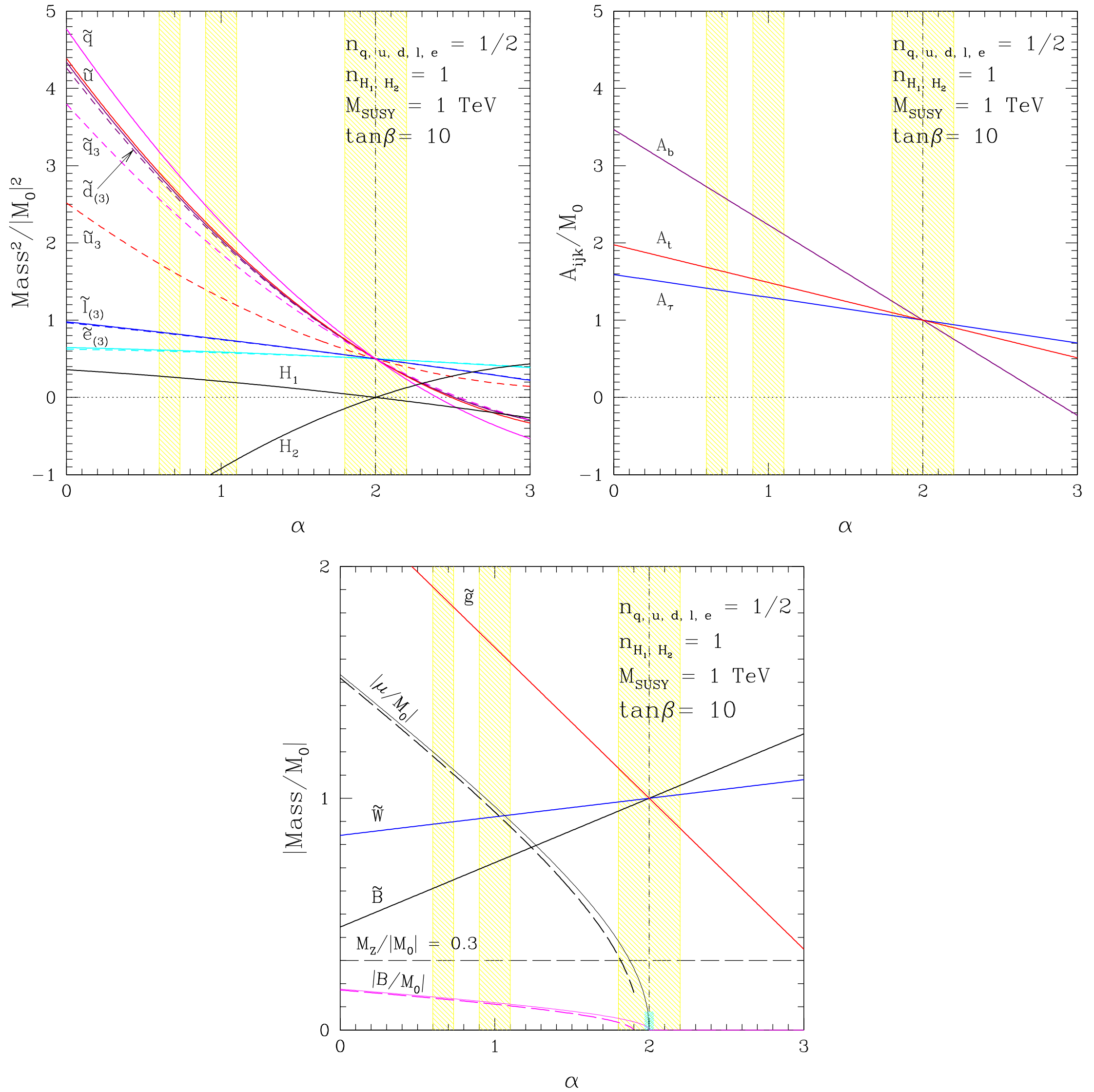

Figure 4: A model for unified soft parameters at M_{SUSY} for $\alpha = 2$. We choose $n_{q,u,d,l,e} = 1/2$ and $n_{H_1,H_2} = 1$ to obtain $a_{ijk} = c_i + c_j + c_k = 1$. A little hierarchy between the Higgs boson masses and the sparticle masses for $\alpha = 2$ is protected at 1-loop level. Note that the results on μ/M_0 and B/M_0 at $\alpha \approx 2$ have an uncertainty of $\mathcal{O}(10^{-1})$ due to the threshold corrections of $\mathcal{O}(M_0^2/8\pi^2)$ to the Higgs mass-squares at M_{GUT}.

[3] V. S. Kaplunovsky and J. Louis, Phys. Lett. B **306**, 269 (1993) [arXiv:hep-th/9303040]; A. Brignole, L. E. Ibanez and C. Munoz, Nucl. Phys. B **422**, 125 (1994) [Erratum-ibid. B **436**, 747 (1995)] [arXiv:hep-ph/9308271].

[4] L. Randall and R. Sundrum, Nucl. Phys. B **557**, 79 (1999) [arXiv:hep-th/9810155]; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [arXiv:hep-ph/9810442]; J. A. Bagger, T. Moroi and E. Poppitz, JHEP **0004**, 009 (2000) [arXiv:hep-th/9911029]; P. Binetruy, M. K. Gaillard and B. D. Nelson, Nucl. Phys. B **604**,

32 (2001) [arXiv:hep-ph/0011081].

[5] K. Choi, J. E. Kim and H. P. Nilles, Phys. Rev. Lett. **73**, 1758 (1994) [arXiv:hep-ph/9404311].

[6] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, Phys. Rev. D **68**, 046005 (2003) [arXiv:hep-th/0301240].

[7] J. P. Derendinger, L. E. Ibanez and H. P. Nilles, Phys. Lett. B **155**, 65 (1985); M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. B **156**, 55 (1985); J. P. Derendinger, L. E. Ibanez and H. P. Nilles, Nucl. Phys. B **267**, 365 (1986); H. P. Nilles, arXiv:hep-th/0402022.

[8] K. Choi, A. Falkowski, H. P. Nilles, M. Olechowski and S. Pokorski, JHEP **0411**, 076 (2004) [arXiv:hep-th/0411066].

[9] K. Choi, A. Falkowski, H. P. Nilles and M. Olechowski, Nucl. Phys. **B718**, 113 (2005) [arXiv:hep-th/0503216].

[10] T. Gherghetta, G. F. Giudice, J. D. Wells, Nucl. Phys. **B559**, 27 (1999) [arXiv:hep-ph/9904378]; J. L. Feng and T. Moroi, Phys. Rev. **D61**, 095004 (2000) [arXiv:hep-ph/9907319].

[11] M. A. Luty and R. Sundrum, Phys. Rev. **D62**, 035008 (2000) [arXiv:hep-th/9910202]; Z. Chacko and M. A. Luty, JHEP **0105**, 067 (2001) [arXiv:hep-ph/0008103].

[12] K. Choi, Phys. Rev. Lett. **72**, 1592 (1994) [arXiv:hep-ph/9311352].

[13] L. E. Ibanez, C. Munoz and S. Rigolin, Nucl. Phys. **B553**, 43 (1999) [arXiv:hep-ph/9812397].

[14] L. E. Ibanez, arXiv:hep-ph/0408064; B. C. Allanach, A. Brignole and L. E. Ibanez, arXiv:hep-ph/0502151.

[15] S. B. Giddings, S. Kachru and J. Polchinski, Phys. Rev. D **66**, 106006 (2002) [arXiv:hep-th/0105097].

[16] S. Samuel and J. Wess, Nucl. Phys. B **221**, 153 (1983); J. Bagger and A. Galperin, Phys. Rev. D **55** (1997) 1091 [arXiv:hep-th/9608177]; M. Klein, Phys. Rev. **D66**, 055009 (2002) [arXiv:hep-th/0205300]; C. P. Burgess, E. Filotas, M. Klein and F. Quevedo, JHEP **0310**, 041 (2003) [arXiv:hep-th/0209190].

[17] C. P. Burgess, R. Kallosh and F. Quevedo, JHEP **0310**, 056 (2003) [arXiv:hep-th/0309187].

[18] S. Kachru, R. Kallosh, A. Linde, J. Maldacena, L. McAllister and S. P. Trivedi, JCAP **0310**, 013 (2003) [arXiv:hep-th/0308055]

[19] N. V. Krasnikov, Phys. Lett. B **193**, 37 (1987); J. A. Casas, Z. Lalak, C.Munoz and G. G. Ross, Nucl. Phys. B **347**, (1990) 243; T. R. Taylor, Phys. Lett. B **252**, (1990) 59; B. de Carlos, J. A. Casas and C. Munoz, Nucl. Phys. B **399** (1993) 623 [arXiv:hep-th/9204012].

[20] G. F. Giudice and A. Masiero, Phys. Lett. **B206**, 480 (1988).

[21] E. J. Chun, J. E. Kim and H. P. Nilles Nucl. Phys. **B370**, 105 (1992).

[22] J. R. Ellis, J. F. Gunion, H. E. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D **39**, 844 (1989).

[23] D. Marti and A. Pomarol, Phys. Rev. **D64**, 105025 (2001) [arXiv:hep-th/0106256].

[24] H. Abe, K. Choi, K. S. Jeong and K. Okumura, JHEP **0409**, 015 (2004) [arXiv:hep-ph/0407005].

[25] P. Binetruy, G. Dvali, R. Kallosh and A. Van Proeyen, Class. Quant. Grav. **21**, 3137 (2004) [arXiv:hep-th/0402046]; H. Abe and K. Choi, JHEP **0412**, 069 (2004) [arXiv:hep-th/0412174].

[26] K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **68**, 927 (1982) [Erratum-ibid. **70**, 330 (1983)].

[27] L. E. Ibanez and G. G. Ross, Phys. Lett. B **110**, 215 (1982); J. R. Ellis, D. V. Nanopoulos and K. Tamvakis, Phys. Lett. B **121**, 123 (1983); L. Alvarez-Gaume, J. Polchinski and M. B. Wise, Nucl. Phys. B **221**, 495 (1983).

[28] S. Mizuta and M. Yamaguchi, Phys. Lett. B **298**, 120 (1993) [arXiv:hep-ph/9208251]; J. L. Feng, K. T. Matchev and F. Wilczek, Phys. Lett. B **482**, 388 (2000) [arXiv:hep-ph/0004043]; P. Gondolo and J. Edsjo, arXiv:hep-ph/9804459; T. Falk, arXiv:hep-ph/9807481.

[29] R. Barbieri and G. F. Giudice, Nucl. Phys. B **306**, 63 (1988);

[30] S. Eidelman *et al.* [Particle Data Group], Phys. Lett. B **592**, 1 (2004).

[31] J. A. Casas, J. R. Espinosa and I. Hidalgo, JHEP **0401**, 008 (2004) [arXiv:hep-ph/0310137].

[32] K. Kohri, M. Yamaguchi and J. Yokoyama, arXiv:hep-ph/0502211.